

High-scale validity of a two Higgs doublet scenario: metastability included.Nabarun Chakrabarty^{†1} and Biswarup Mukhopadhyaya^{★2}[†]*Regional Centre for Accelerator-based Particle Physics**Harish-Chandra Research Institute**Chhatnag Road, Jhansi, Allahabad - 211 019, India***Abstract**

We make an attempt to identify regions in a Type II Two-Higgs Doublet Model, which correspond to a metastable electroweak vacuum with lifetime larger than the age of the universe. We analyse scenarios which retain perturbative unitarity up to Grand unification and Planck scales. Each point in the parameter space is restricted using Data from the Large Hadron Collider (LHC) as well as flavor and precision electroweak constraints. We find that substantial regions of the parameter space are thus identified as corresponding to metastability, which compliment the allowed regions for absolute stability, for top quark mass at the high as well as low end of its currently allowed range. Thus, a two-Higgs doublet scenario with the electroweak vacuum, either stable or metastable, can sail through all the way up to the Planck scale without facing any contradictions.

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1 Introduction

With the observation of a scalar resonance around 125 GeV at the LHC [1, 2], and hence its identification with a Higgs boson, the particle spectrum of the Standard Model (SM) appears to be complete. However, issues ranging from the existence of Dark Matter (DM) to the pattern of neutrino mass continue to suggest physics beyond the SM. While the quest for such new physics remains on, a rather pertinent question to ask is whether the SM by itself can ensure vacuum stability at scales above that of electroweak symmetry breaking. This is because the Higgs quartic coupling evolving via SM interaction alone tends to turn negative in between the Electroweak (EW) and Planck scales, thereby making the scalar potential unbounded from below. This exact location of this instability scale crucially depends on the pole masses of the top quark and the Higgs. A recent next-to-next-to-leading order (NNLO) study [3, 4] finds that absolute stability up to the Planck scale requires [3]

$$M_h[\text{GeV}] > 129.4 + 1.4\left(\frac{M_t[\text{GeV}] - 173.1}{0.7}\right) - 0.5\left(\frac{\alpha_s(M_Z) - 0.1184}{0.0007}\right) \pm 1.0_{\text{th}} \quad (1.1)$$

The updated measurements of the Higgs and top quark masses [REF] hint towards a *metastable* vacuum scenario. Such a scenario leads to an additional minimum in the potential, which is deeper than the electroweak vacuum [5]. However the lifetime of the latter is less than the age of the universe, thus enabling the present day vacuum to be consistent with the well tested electroweak theory.

In general, vacuum instability can be alleviated by introducing additional bosonic degrees of freedom, which can offset the downward evolution of the quartic coupling of the SM. Such a possibility has indeed been explored in the context of various non-minimal Higgs sectors. One example would be the case of the celebrated Two-Higgs Doublet Models (2HDMs). Different types of 2HDM offer interesting phenomenology at present and future colliders and, are consistent with flavor physics constraints, and are part and parcel of supersymmetric theories. In one of our earlier works, we showed in context of a Type-II 2HDM that the EW vacuum can be rendered stable till the Planck scale even for a top pole mass at the high end of the allowed band [6]. Moreover, this can be achieved without running into conflict with perturbativity or unitarity at high scales.

A 2HDM is set has one more vital feature. The Yukawa couplings of the SM fermions can be different compared to the SM values, in a 2HDM. The role of the coupling for the same M_t can thus produce a different effect to the evolution to the quartics. A stable vacuum till the Planck scale is achieved with specific combinations of the boundary conditions, i.e, given in terms of the model parameters. This raises the question, *is there a possible metastable vacuum in a 2HDM? Can such a balance between the bosonic and fermionic effects be struck that indeed leads to an additional minimum of the scalar potential, while prolonging the lifetime of the EW vacuum to a safe level? This is the question we precisely seek to answer in this work.*

Metastability in context of certain non-minimal Higgs scenarii has been looked into in the recent times and along with the attempts of embedding successful cold dark matter candidates in the parameter space that leads to a stable/metastable vacuum. These studies however confront *inert* scalars that do not mix with the SM Higgs owing to some discrete symmetry, and moreover, the Higgs-top coupling also remains unchanged with respect to the SM. The implications could be different in a 2HDM where EWSB is triggered when both the doublets receive vacuum expectation values (vev). Two possibilities open thus up: (a) The scalar potential could furnish additional neutral minima around the TeV scale ballpark in the slice spanned

by the neutral fields in the two doublets, and, (b) Additional minimum can appear when the scalar potential is improved by Renormalisation Group (RG) effects. In the process of looking for such possibilities, we take into account miscellaneous constraints coming from Higgs signal strengths, flavor issues and electroweak precision observables.

The paper has the following plan. In Sec. 2, we review the salient features of the 2HDMs. Sec. 3 is dedicated to a discussion on how a metastable vacuum can arise, and on the completion of its lifetime. We also present an outline of the tunnelling probability computation in the same. Sec. 4 presents an overall strategy on how to look for a metastable vacua, and, also an account of the various experimental and theoretical constraints taken while doing so. The numerical results are highlighted in Sec. 5 and finally the study is concluded in Sec. 6.

2 Model features.

2.1 Scalar potential

In the present work, we consider the most general renormalizable scalar potential for two doublets Φ_1 and Φ_2 , each having hypercharge (+1),

$$\begin{aligned}
V(\Phi_1, \Phi_2) = & m_{11}^2 \Phi_1^\dagger \Phi_1 + m_{22}^2 \Phi_2^\dagger \Phi_2 - m_{12}^2 (\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1) + \frac{\lambda_1}{2} (\Phi_1^\dagger \Phi_1)^2 + \frac{\lambda_2}{2} (\Phi_2^\dagger \Phi_2)^2 \\
& + \lambda_3 \Phi_1^\dagger \Phi_1 \Phi_2^\dagger \Phi_2 + \lambda_4 \Phi_1^\dagger \Phi_2 \Phi_2^\dagger \Phi_1 + \frac{\lambda_5}{2} \left[(\Phi_1^\dagger \Phi_2)^2 + (\Phi_2^\dagger \Phi_1)^2 \right] \\
& + \lambda_6 \Phi_1^\dagger \Phi_1 (\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1) + \lambda_7 \Phi_2^\dagger \Phi_2 (\Phi_1^\dagger \Phi_2 + \Phi_2^\dagger \Phi_1). \tag{2.1}
\end{aligned}$$

This scenario in general has the possibility of CP-violation in the scalar sector, through the phases in m_{12} , λ_5 , λ_6 and λ_7 . We choose m_{12} to be real here, moreover the terms proportional to λ_6 and λ_7 have been neglected in the present study.

In a general two-Higgs-doublet model (2HDM), a particular fermion can couple to both Φ_1 and Φ_2 . However this would lead to the flavor changing neutral currents (FCNC) at the tree level¹. One way to avoid such FCNC is to impose a \mathbb{Z}_2 symmetry, such as one that demands invariance under $\Phi_1 \rightarrow -\Phi_1$ and $\Phi_2 \rightarrow \Phi_2$. This type of symmetry puts restrictions on the scalar potential. The \mathbb{Z}_2 symmetry is *exact* as long as m_{12} , λ_6 and λ_7 vanish, when the scalar sector also becomes CP-conserving. The symmetry is said to be broken *softly* if it is violated in the quadratic terms only, i.e., in the limit where λ_6 and λ_7 vanish but m_{12} does not. Finally, a *hard* breaking of the \mathbb{Z}_2 symmetry is realized when it is broken by the quartic terms as well. Thus in this case, m_{12} , λ_6 and λ_7 all are non-vanishing in general.

We focus on a particular scheme of coupling fermions to the doublets. This scheme is known in the literature as the *Type-II* 2HDM, where the down type quarks and the charged leptons couple to Φ_1 and the up type quarks, to Φ_2 [18]. This is ensured through the discrete symmetry $\Phi_1 \rightarrow -\Phi_1$ and $\psi_R^i \rightarrow -\psi_R^i$, where ψ is charged leptons or down type quarks and i represents the generation index.

¹In context of a typical flavour changing scenario, it has been shown recently that the FCNCs are stable under RG evolution mostly.

Minimization of the scalar potential in Eqn. 3.6 leads to

$$\langle \Phi_1 \rangle = \begin{pmatrix} 0 \\ \frac{v_1}{\sqrt{2}} \end{pmatrix}, \quad \langle \Phi_2 \rangle = \begin{pmatrix} 0 \\ \frac{v_2}{\sqrt{2}} \end{pmatrix}, \quad (2.2)$$

where the vacuum expectation values (vev) are often expressed in terms of the M_Z and the ratio

$$\tan \beta = \frac{v_2}{v_1}. \quad (2.3)$$

We parametrise the doublets in the following fashion,

$$\Phi_i = \frac{1}{\sqrt{2}} \begin{pmatrix} \sqrt{2} w_i^+ \\ v_i + h_i + i z_i \end{pmatrix} \text{ for } i = 1, 2. \quad (2.4)$$

Since the basis used in $V(\Phi_1, \Phi_2)$ allows mixing between the two doublets, the physical states are obtained by diagonalising the charged and neutral scalar mass matrices. There are then altogether eight mass eigenstates, three of which become the longitudinal components of the W^\pm and Z gauge bosons. Of the remaining five, there is a mutually conjugate pair of charged scalars (H^\pm), two neutral CP even scalars (H, h) and a neutral pseudoscalar (A), given there is no CP-violation. Otherwise, a further mixing between (H, h) and A becomes unavoidable. The compositions of the mass eigenstates H and h indeed depend on the mixing angle α .

The quartic couplings are conveniently expressed in terms of the physical masses and mixing angles as,

$$\lambda_1 = \frac{1}{v^2 c_\beta^2} \left(c_\alpha^2 m_H^2 + v^2 s_\alpha^2 m_h^2 - m_{12}^2 \frac{s_\beta}{c_\beta} \right) \quad (2.5a)$$

$$\lambda_2 = \frac{1}{v^2 s_\beta^2} \left(s_\alpha^2 m_H^2 + v^2 c_\alpha^2 m_h^2 - m_{12}^2 \frac{c_\beta}{s_\beta} \right) \quad (2.5b)$$

$$\lambda_3 = \frac{2m_{H+}^2}{v^2} + \frac{s_{2\alpha}}{v^2 s_{2\beta}} (m_H^2 - m_h^2) - \frac{m_{12}^2}{v^2 s_\beta c_\beta} \quad (2.5c)$$

$$\lambda_4 = \frac{1}{v^2} (m_A^2 - 2m_{H+}^2) + \frac{m_{12}^2}{v^2 s_\beta c_\beta} \quad (2.5d)$$

$$\lambda_5 = \frac{m_{12}^2}{v^2 s_\beta c_\beta} - \frac{m_A^2}{v^2} \quad (2.5e)$$

3 The computation of tunneling probability.

The existence of a large number of scalar degrees of freedom makes the vacuum landscape of a 2HDM more elaborate and intriguing compared to the SM. Here we are confining ourselves to the situation when the vacuum breaks neither electric charge nor CP. Under such circumstances, the EWSB conditions appear as,

$$m_{11}^2 v_1 = m_{12}^2 v_2 - \frac{1}{2} \lambda_1 v_1^3 - \frac{1}{2} (\lambda_3 - \lambda_4 + \lambda_5) v_1 v_2^2 \quad (3.1)$$

$$m_{22}^2 v_2 = m_{12}^2 v_1 - \frac{1}{2} \lambda_2 v_2^3 - \frac{1}{2} (\lambda_3 - \lambda_4 + \lambda_5) v_2 v_1^2 \quad (3.2)$$

It has been reported in [7] that the above conditions can lead to several solutions, and at most two non-degenerate minima. In other words, apart from the EW minimum in which the universe currently resides

($v_1^2 + v_2^2 = 246 \text{ GeV}^2$, named N), there exists another minimum somewhere around ($v_1^2 + v_2^2 \neq 246 \text{ GeV}^2$, named N'). Ref. finds the difference of depths of the tree level scalar potential at the two minima to be,

$$V_{N'} - V_N = \frac{m_{12}^2}{4v_1v_2} \left(1 - \frac{v_1v_2}{v'_1v'_2}\right)^2 (v_1v'_2 - v_2v'_1)^2 \quad (3.3)$$

Thus there exists the tantalizing possibility that the 2HDM offers such parameter points for which a neighbouring vacuum could actually be deeper than the one which corresponds to the observed W - and Z -boson masses. The EW minimum then loses its status as the global minimum and has been termed the *panic vacuum* in [7]. In those cases, computing the lifetime of tunnelling to the non-EW minimum from the EW one becomes the pertinent task. If the tunnelling lifetime turn out to be higher than the age of the universe, the non-EW minimum cannot be ruled out. However, thanks to the data from the LHC on Higgs signal strengths, the model points admitting $V_{N'} - V_N < 0$ are more or less ruled out [7, 8]. However, a new landscape of vacua can still open up if one investigates the *renormalisation-group improved* effective potential in place of the bare tree-level one. In the context of the SM, it can be understood as follows: The SM quartic coupling turns negative at some energy scale 10^{8-11} GeV (exact location of the scale depends on the choice of the initial conditions), after which it again starts rising owing to the bosonic effects counterbalancing the negative top-Yukawa drag. The fallout of this is the emergence of a new minimum beyond the scale where the quartic coupling first becomes negative. It should be noted that the direction of the EW vacuum uniquely decides the direction in which the high-scale vacuum is formed.

In a 2HDM, on the other hand, one has to handle the additional complication of having a higher number of field directions. In addition, the effects of the various interaction terms make it imperative to incorporate the effects of radiative corrections induced by the 2HDM. Therefore, we choose to analyse the one-loop corrected effective potential in place of the tree level potential. One thus writes

$$V_{eff}(h_1, h_2) = V_{tree}(h_1, h_2) + V_{1loop}(h_1, h_2) \quad (3.4)$$

Here, $V_{tree}(h_1, h_2)$ and $V_{1loop}(h_1, h_2)$ denote the tree level and 1-loop parts of the effective potential, calculated along the $h_1 - h_2$ subspace.

For example, , the tree level potential reads

$$V_{tree}(h_1, h_2) = \frac{1}{2}m_{11}^2h_1^2 + \frac{1}{2}m_{22}^2h_2^2 + m_{12}^2h_1h_2 + \frac{\lambda_1}{8}h_1^4 + \frac{\lambda_2}{8}h_2^4 + \frac{\lambda_3 + \lambda_4 + \lambda_5}{4}h_1^2h_2^2 \quad (3.5)$$

In the (h_1, h_2) plane, it has the following expression,

$$V_{1loop}(h_1, h_2) = \frac{1}{64\pi^2} \sum_i n_i M_i^4(h_1, h_2) \left[\ln\left(\frac{M_i^4(h_1, h_2)}{\mu^2}\right) - c_i \right] \quad (3.6)$$

Where n_i refers to the number of degrees of freedom for the i th field and c_i are constants whose values depend on the regularization scheme adopted. To list the constants explicitly, $n_W = 6$, $n_Z = 3$, $n_t = -12$, $n_h = 1$; and $c_W = \frac{5}{6}$, $c_Z = \frac{5}{6}$, $c_t = \frac{3}{2}$, $c_h = \frac{3}{2}$. Moreover μ refers to the renormalization scale emerging as an artifact of dimensional regularization. $M_i^2(h_1, h_2)$ represent the scale dependent mass squared.

Studies on the high-scale validity of a 2HDM in the past were mostly confined to investigating absolute stability [9–12]. Some studies connecting higher dimensional operators to Higgs metastability have occurred in the past [13, 14]. The main theme of this work is to investigate possible high scale vacua in the context of 2HDM using the prescription suggested by Coleman. $V_{eff}(h_1, h_2)$ depends on two variables, and hence,

determining a classical solution interpolating the two vacua, even numerically, becomes an extremely challenging task. Furthermore, a generic classical path may not qualify as a "bounce" [15, 16], i.e, it might not pass through the top of a barrier separating two vacua. Coleman's prescription does not apply in such a case. However, one can always choose to look for additional minima along a particular ray in the $h_1 - h_2$ plane. Under this approximation, the effective potential is reduced to a function of a single variable again (that particular linear combination of h_1 and h_2). In models such as Type I or Type II 2HDM, the Z_2 symmetry of the Yukawa interactions implies that the top quark always couples to Φ_2 . Thus it is the coupling λ_2 that experiences the maximum downward pull due to the Yukawa interactions and can consequently turn negative at high scales in spite of starting with a positive value at the input scale. It therefore makes sense to look for additional minima in the h_2 direction only. This approach is similar to what [17] opts in context of an inert doublet model.

We study the behaviour of the $V_{eff}(h_1, h_2)$ in the limit where $h_1 \simeq v$ and $h_2 \gg h_1, m_{12}$. In this limit, the squared masses have the following simplified expressions:

$$m_{H_1}^2(h_2) \simeq \frac{1}{2}(\lambda_3 + \lambda_4 + \lambda_5)h_2^2 \quad (3.7)$$

$$m_{H_2}^2(h_2) \simeq \frac{3}{2}\lambda_2 h_2^2 \quad (3.8)$$

$$m_{A_1}^2(h_2) \simeq \frac{1}{2}(\lambda_3 + \lambda_4 - \lambda_5)h_2^2 \quad (3.9)$$

$$m_{A_2}^2(h_2) \simeq \frac{1}{2}\lambda_2 h_2^2 \quad (3.10)$$

$$m_{H_1^+}^2(h_2) \simeq \frac{1}{2}\lambda_3 h_2^2 \quad (3.11)$$

$$m_{H_2^+}^2(h_2) \simeq \frac{1}{2}\lambda_2 h_2^2 \quad (3.12)$$

$$m_t^2(h_2) \simeq \frac{1}{2}y_t^2 h_2^2 \quad (3.13)$$

$$m_W^2(h_2) \simeq \frac{1}{4}g^2 h_2^2 \quad (3.14)$$

$$m_Z^2(h_2) \simeq \frac{1}{4}(g^2 + g'^2)h_2^2 \quad (3.15)$$

All running couplings are evaluated at the scale $\mu \simeq h_2$. Under all these approximations, the real part of one-loop corrected potential takes the form,

$$V_{eff}(h_2) \simeq \frac{\lambda_2^{eff}}{8}h_2^4 \quad (3.16)$$

$$\text{where,} \quad (3.17)$$

$$\begin{aligned} \lambda_2^{eff}(h_2) \simeq & \lambda_2(h_2) + \frac{1}{64\pi^2} \left[2(\lambda_3 + \lambda_4 + \lambda_5)^2 \left(\ln \frac{\lambda_3 + \lambda_4 + \lambda_5}{2} - \frac{3}{2} \right) + 18\lambda_2^2 \left(\ln \frac{3\lambda_2}{2} - \frac{3}{2} \right) \right. \\ & 2(\lambda_3 + \lambda_4 - \lambda_5)^2 \left(\ln \frac{\lambda_3 + \lambda_4 - \lambda_5}{2} - \frac{3}{2} \right) + 2\lambda_3^2 \left(\ln \frac{\lambda_3}{2} - \frac{3}{2} \right) + 6\lambda_2^2 \left(\ln \frac{\lambda_3}{2} - \frac{3}{2} \right) \\ & \left. 3g^4 \left(\ln \frac{g^2}{4} - \frac{5}{6} \right) + \frac{3}{2}(g^2 + g'^2)^2 \left(\ln \frac{g^2 + g'^2}{4} - \frac{5}{6} \right) - 24y_t^4 \left(\ln \frac{y_t^2}{2} - \frac{3}{2} \right) \right] \end{aligned}$$

Where the term in square brackets refers to the finite correction generated by the Coleman Weinberg mechanism. We find that highly sub-dominant in our calculations.

The probability of tunneling to the deeper vacuum is given by,

$$p = T_U^4 \mu^4 e^{-\frac{8\pi^2}{3|\lambda_2^{eff}|}} \quad (3.18)$$

Here μ refers to the scale where the probability is maximized, and, it turns out that $\frac{d\lambda_2}{d\log(Q)} = 0$ at $Q = \mu$. Using $T_U \simeq 10^{10}$ yr and requiring that the vacuum tunnelling lifetime is always higher than the lifetime on the universe tantamounts to having the following condition [5],

$$\lambda_2^{eff}(h_2) \geq \frac{-0.065}{1 - 0.01 \ln(a)} \quad (3.19)$$

where $a = \frac{v}{\mu}$. It may be noted that we have accepted λ_2 turning negative in the h_2 direction as the sole condition for the loss of stability of the EW vacuum. There is in general an extended set of conditions for stability in a 2HDM [18]. However, one can easily verify that the remaining conditions for stability in a 2HDM are violated, if at all, at low scale itself. Such violation, on the other hand, leads to the disappearance of the EW minimum as a whole. This cannot be a situation appropriate for metastability, and therefore the conditions other than $\lambda_2 < 0$ need not be used as signs for loss of stability.

4 A metastable vacuum and the 2HDM parameter space.

4.1 Analysis strategy.

As has been already in the previous section a look-out for an additional vacuum at high scales requires one to study the evolution of the various interaction strengths under renormalization group equations. The values of the quartic couplings at the electroweak scale are, of course, connected with the masses and mixing angles in the scalar sector. A careful measurement of the signal strengths of the 125 GeV at the LHC has revealed that the resonance has couplings strikingly similar to the SM ones. These observations have their ramifications on the 2HDM parameter space. Thus together with the requirement of having $m_h \simeq 125$ GeV, we also arrange for $\beta - \alpha \simeq \frac{\pi}{2}$, in order to comply with these results from the LHC ($\beta - \alpha = \frac{\pi}{2}$ is the well known *alignment limit* [19] in a 2HDM, in which the couplings of h to fermions and gauge bosons become exactly equal to the SM ones). We choose to describe a 2HDM parameter point in terms of the parameters $(\tan \beta, m_h, m_H, m_A, m_{H^\pm}, \alpha)$ basis. There are additional constraints to satisfy as outlined in the next few subsections,

4.1.1 Perturbativity, unitarity and vacuum stability

For the 2HDM to remain a perturbative quantum field theory at a given energy scale, one must impose the conditions $|\lambda_i| \leq 4\pi$ ($i = 1, \dots, 7$) and $|y_i| \leq \sqrt{4\pi}$ ($i = t, b, \tau$) at that scale². This translates into upper bounds on the running couplings at low as well as high scales.

A more sophisticated version of such bounds comes from the requirement of partial wave unitarity in longitudinal gauge boson scattering. The $2 \rightarrow 2$ amplitude matrix corresponding to scattering of the longitudinal components of the gauge bosons can be mapped to a corresponding matrix for the scattering of the goldstone bosons [20, 21]. The theory respects unitarity if each eigenvalue of the aforementioned amplitude matrix does not exceed 8π .

$$a_{\pm} = \frac{3}{2}(\lambda_1 + \lambda_2) \pm \sqrt{\frac{9}{4}(\lambda_1 - \lambda_2)^2 + (2\lambda_3 + \lambda_4)^2}, \quad (4.1a)$$

²The conditions are slightly different for the two types of couplings. The reason becomes clear if we note that the perturbative expansion parameter for $2 \rightarrow 2$ processes driven by the quartic couplings is λ_i . The corresponding parameter for Yukawa-driven scattering processes is $|y_i|^2$

$$b_{\pm} = \frac{1}{2}(\lambda_1 + \lambda_2) \pm \sqrt{\frac{1}{4}(\lambda_1 - \lambda_2)^2 + \lambda_4^2}, \quad (4.1b)$$

$$c_{\pm} = d_{\pm} = \frac{1}{2}(\lambda_1 + \lambda_2) \pm \sqrt{\frac{1}{4}(\lambda_1 - \lambda_2)^2 + \lambda_5^2}, \quad (4.1c)$$

$$e_1 = (\lambda_3 + 2\lambda_4 - 3\lambda_5), \quad (4.1d)$$

$$e_2 = (\lambda_3 - \lambda_5), \quad (4.1e)$$

$$f_1 = f_2 = (\lambda_3 + \lambda_4), \quad (4.1f)$$

$$f_+ = (\lambda_3 + 2\lambda_4 + 3\lambda_5), \quad (4.1g)$$

$$f_- = (\lambda_3 + \lambda_5). \quad (4.1h)$$

When the quartic part of the scalar potential preserves CP and \mathbb{Z}_2 symmetries, the aforementioned eigenvalues are discussed in [20–22].

The condition to be discussed next is that of vacuum stability. For the scalar potential of a theory to be stable, it must be bounded from below in all possible directions. This condition is threatened if the quartic part of the scalar potential, which is responsible for its behaviour at large field values, turns negative. Avoiding such a possibility till any given scale ensures stability of the vacuum up to that scale. Vacuum stability in context of a 2HDM has been discussed in detail in [7]

Demanding high-scale positivity of the 2HDM potential along various directions in the field space leads to the following conditions on the scalar potential. [18, 23],

$$\text{vsc1} : \lambda_1 > 0 \quad (4.2a)$$

$$\text{vsc2} : \lambda_2 > 0 \quad (4.2b)$$

$$\text{vsc3} : \lambda_3 + \sqrt{\lambda_1 \lambda_2} > 0 \quad (4.2c)$$

$$\text{vsc4} : \lambda_3 + \lambda_4 - |\lambda_5| + \sqrt{\lambda_1 \lambda_2} > 0 \quad (4.2d)$$

When a second vacuum arises, we go for a test of metastability of the electroweak vacuum using vsc2 alone. The reason for this has already been discussed at the end of section 3.

4.1.2 Oblique parameters and flavor constraints.

A 2HDM contributes to the electroweak precision observables through the participation of the additional scalars in loops. For example, the oblique S , T and U parameters receive contributions ΔS , ΔT and ΔU respectively from the 2HDM. The most constraining amongst these is ΔT which reads,

$$\begin{aligned} \Delta T = & F(m_{H^+}^2, m_H^2) + F(m_{H^+}^2, m_A^2) + c_{\beta-\alpha}^2 F(m_{H^+}^2, m_h^2) + s_{\beta-\alpha}^2 F(m_{H^+}^2, m_H^2) - F(m_H^2, m_A^2) \\ & - c_{\beta-\alpha}^2 F(m_h^2, m_A^2) - s_{\beta-\alpha}^2 F(m_H^2, m_A^2) + 3c_{\beta-\alpha}^2 (F(m_Z^2, m_H^2) - F(m_W^2, m_H^2)) \\ & - 3c_{\beta-\alpha}^2 (F(m_Z^2, m_h^2) - F(m_W^2, m_h^2)) \end{aligned} \quad (4.3a)$$

Where,

$$F(m_1^2, m_2^2) \equiv \begin{cases} \frac{m_1^2 + m_2^2}{2} - \frac{m_1^2 m_2^2}{m_1^2 - m_2^2} \ln \frac{m_1^2}{m_2^2} & ; \quad m_1^2 \neq m_2^2, \\ 0 & ; \quad m_1^2 = m_2^2. \end{cases} \quad (4.4)$$

We have filtered the model points through the constraint $\Delta T = 0.05 \pm 0.12$ following [24]. On the restrictions coming from the flavor physics side, measurement of the $b \rightarrow s\gamma$ leads to $m_{H^+} \geq 315$ GeV in case of the Type-II 2HDM [25]. In case of Type-I, there is no such lower bound. The constraint $m_{H^+} \geq 80$ GeV originating from direct searches however still persists.

4.1.3 $h \rightarrow \gamma\gamma$ decay width

For a near-alignment case, the Higgs signal strengths to The partial decay width of the SM-like Higgs to a pair of photons in this case has the expression [26],

$$\Gamma^{2HDM}(h \rightarrow \gamma\gamma) = \frac{\alpha^2 g^2}{2^{10} \pi^3} \frac{m_h^3}{M_W^2} \left| \sin(\beta - \alpha) F_W + \left(\frac{\cos \alpha}{\sin \beta} \right) \frac{4}{3} F_t + \kappa F_{H^+} \right|^2, \quad (4.5)$$

The functions F_W , F_t and F_{i+} encapsulate the effects of a W-boson, a t-quark and a charged scalar running in the loop and shall be defined as,

$$F_W = 2 + 3\tau_W + 3\tau_W(2 - \tau_W)f(\tau_W), \quad (4.6a)$$

$$F_t = -2\tau_t[1 + (1 - \tau_t)f(\tau_t)], \quad (4.6b)$$

$$F_{H^+} = -\tau_{i+}[1 - \tau_{i+}f(\tau_{i+})]. \quad (4.6c)$$

$$f(\tau) = \left[\sin^{-1} \left(\sqrt{1/\tau} \right) \right]^2. \quad (4.7)$$

$$\text{with, } \tau = \frac{4m_a^2}{m_h^2} \quad (4.8)$$

Here, $a = t, W$ and H^+ .

We assume h is dominantly produced through gluon fusion. In such a case, the signal strength for the diphoton final state is approximately given by,

$$\mu_{\gamma\gamma} = \frac{\Gamma^{2HDM}(h \rightarrow \gamma\gamma)}{\Gamma^{SM}(h \rightarrow \gamma\gamma)} \quad (4.9)$$

In order to respect the 2σ bound on $\mu_{\gamma\gamma}$, from a combined measurement of ATLAS and CMS, we discard model points that violate $\mu_{\gamma\gamma} \in [1.04, 1.37]$ [27]

5 Results and discussions.

Model points that successfully negotiate all the aforesaid constraints are allowed to evolve under RG, till some scale Λ (say). Λ is essentially identified as the scale where perturbative unitarity is destroyed, and can be interpreted as the scale up to which no physics over and above the extended Higgs sector is required. If there is an additional, lower vacuum before Λ , the time scale for tunneling from the EW vacua to the new one must therefore be larger than the age of the universe. It is intuitively expected that higher is Λ , tighter becomes the parameter space that is allowed at the electroweak scale. This is indeed confirmed by the findings reported in REF. Of course the points leading to a metastable EW vacuum are identified through a detailed scan of the parameter space. However, the fate of a particular model point at high scales is sensitive to the value of the top quark mass taken. With this in view, we propose the benchmarks listed in Table 1.

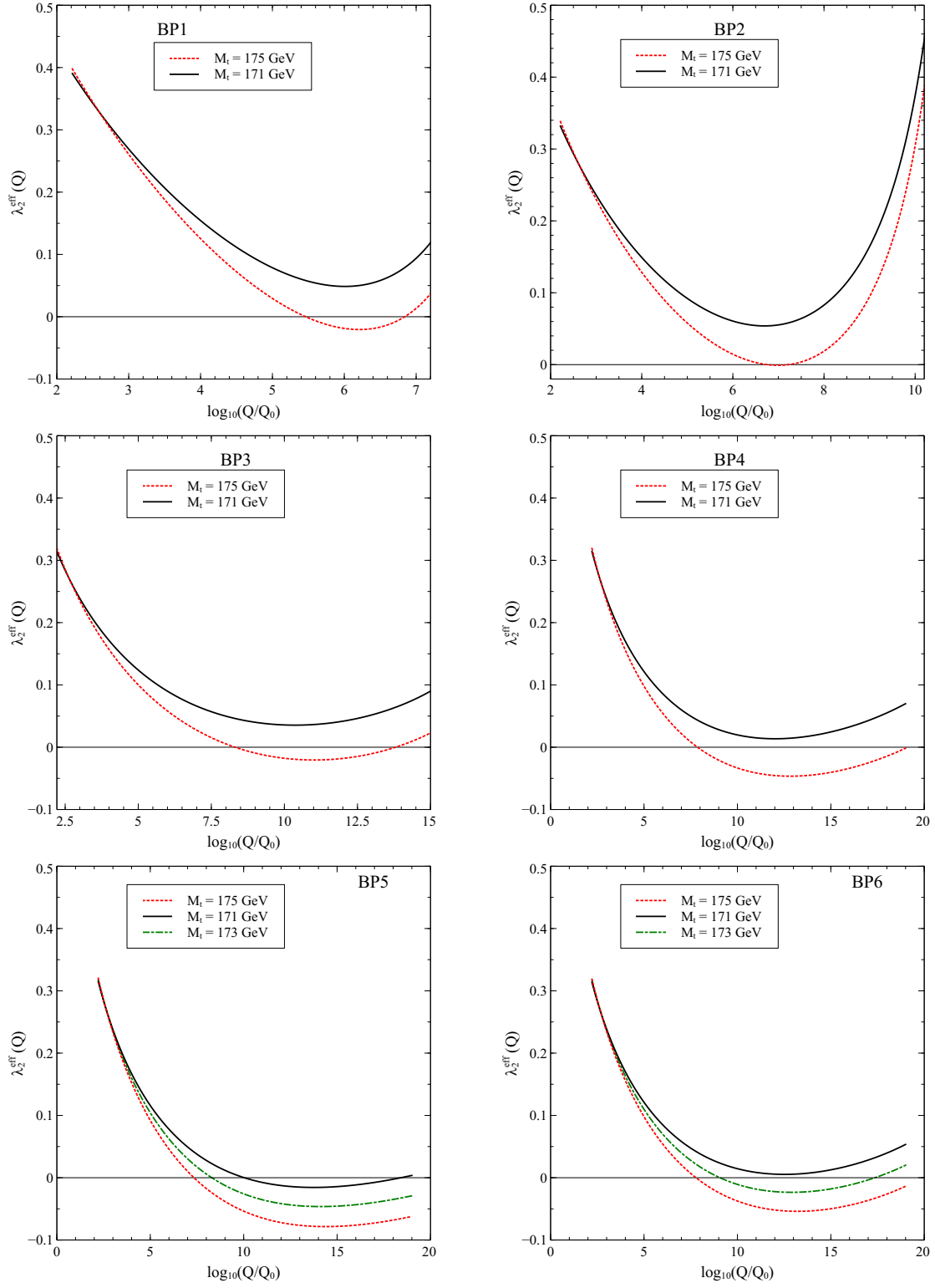


Figure 1: RG evolution of λ_2^{eff} for the benchmarks listed in Table 1, for more than one value of M_t . The color coding is explained in the legends.

Benchmark	$\tan\beta$	$m_H(\text{GeV})$	$m_A(\text{GeV})$	$m_{H^+}(\text{GeV})$	$m_{12}(\text{GeV})$	Perturbative till
BP1	1.78	354	380	341	222	$\sim 10^7$ GeV
BP2	2.50	489	506	486	286	$\sim 10^{11}$ GeV
BP3	7.28	320	297	324	117	$\sim 10^{16}$ GeV
BP4	8.28	500	500	500	172	$\sim 10^{19}$ GeV
BP5	6.90	501	499	500	189	$\sim 10^{19}$ GeV
BP6	10.94	1499	1500	1498	451	$\sim 10^{19}$ GeV

Table 1: Benchmark points chosen to illustrate the behaviour under RGE. Λ denotes the maximum extrapolation scale up to which perturbativity remains intact.

In each case, we plot the evolution of λ_2^{eff} in Fig 1. The chosen benchmarks differ in their perturbative behaviour at high scales, although all of them have the common feature that the EW vacuum turns metastable, or even unstable for $M_t = 175$ GeV. In other words, the Type-II 2HDM may turn non-perturbative beyond a scale Λ , even though a vacuum deeper than the EW one might be encountered someplace intermediate between the electroweak scale and Λ . Besides, although it is worth identifying those parameter points that keep the EW vacuum metastable all the way till the GUT or Planck scales, we also include for completeness in the benchmarks, two points where a 2HDM loses its perturbativity at a much lower scale. For instance in BP1, at $M_t = 175$ GeV, λ_2^{eff} turns negative and, $\frac{d\lambda_2^{eff}}{dt} = 0$ occurs around 6.2×10^6 GeV (The scale at which the tunneling probability gets maximized). Using Eqn, one obtains $\lambda_2^{eff} < -0.05$ in this case. An inspection of Fig 1 thus indicates that this particular benchmark leads to metastability. The same parameter point offers absolute stability for $M_t = 171$ GeV though. BP2 describes the same qualitative feature as BP1, albeit it remains perturbative till 10^{11} GeV.

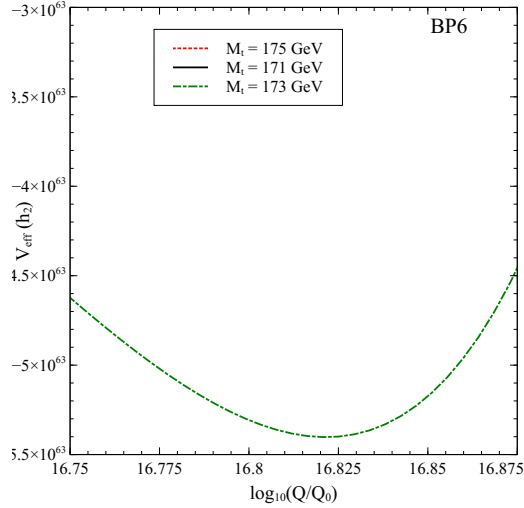


Figure 2: Behaviour of $V_{eff}(h_2)$ in BP6 for $M_t = 173$ GeV.

BP3 is a more conservative benchmark in the sense that, it keeps the 2HDM perturbative till the GUT

scale and also prevents an *unstable* EW vacuum even in the worst case scenario with $M_t = 175$ GeV. We remind the reader that the strength of the top quark Yukawa coupling depends not only on the pole mass, but also on $\tan\beta$. This becomes crucial in deciding the fate of the EW vacuum at high scales. For instance, BP5 experiences a higher t-quark negative pull compared to BP4 owing to a lower value of $\tan\beta$ in BP5, even though the quartic couplings at the input scale are at the same ball-park for the two cases. BP6 is a fine-tuned parameter point that is perturbative till the Planck scale, and for which the EW vacuum is stable, metastable or unstable for $M_t = 171$ GeV, 173 GeV and 175 GeV respectively. For the sake of completeness, we display the behaviour $V_{eff}(h_2)$ for $M_t = 173$ GeV case in Fig 2.

Model points are randomly sampled in the following specified ranges,

$$\begin{aligned} \tan\beta &\in [0.1, 20.0] \\ m_H, m_A, m_{H^+}, m_{12} &\in [0, 1200 \text{ GeV}] \end{aligned} \quad (5.1)$$

A condition forbidding the loss of perturbativity/unitarity at scale Λ is imposed throughout the scan. The following broad features emerge from Fig 3 and Fig. 4. The results are shown for Fig. 3 onwards two representative values of Λ , namely 10^{16} GeV (Fig.3) and $\Lambda = 10^{19}$ GeV (Fig.4).

(i) Perturbativity puts stringent constraints on the splitting amongst the masses. In fact, for $\Lambda = 10^{19}$ GeV the masses are near-degenerate. This effect can be attributed to the fact that for a large mass a splitting, the couplings are already large at the input scale, leading to a blow-up soon after. An individual mass does not get bounded from the above however. Thus, a perturbative theory till high scale automatically respects the T-parameter constraint.

(ii) A smaller $\tan\beta$ for the same M_t implies an enhanced fermionic contribution to the evolution of λ_2 , and hence it favors a metastable vacuum over an absolutely stable one. Consequently, $\tan\beta$ is bounded from below in order to prevent tunnelling to the lower vacuum. Moreover, one would apprehend that the bound by obtained by demanding absolute stability of the EW vacuum to be the stronger than the one obtained when one allows for a metastable scenario. For instance, for $M_t = 171$ GeV, the lower bounds read $\simeq 2.1$ and $\simeq 2.5$ for the two cases.

(iii) The lower bound on $\tan\beta$ of course depends on the choice of M_t . For instance the parameter point parametrised in terms of the masses and $\tan\beta$ indeed shall have different evolution trajectories for two different values of M_t . This is reflected in the upper and lower plots of Fig 3, where one witnesses a tighter lower bound, for both the "stable" as well as "metastable" models. Of course, in this case too, absolute stability yields a stronger bound than metastability. For this value of M_t , any model with $\tan\beta < 2.6$ yields a tunneling lifetime lower than the age of the universe.

(iv) Although the lower bound on $\tan\beta$ should also depend on the Λ chosen, it hardly changes with respect to the 10^{16} GeV value for $\Lambda = 10^{19}$ GeV. Only the number of allowed points shrinks to some extent, other essential features are unchanged.

In the plane of $\tan\beta$ vs masses, it is expected that a particular parameter point responsible for a metastable EW vacuum can always be found in the vicinity of a point that leads to absolute stability, $\tan\beta \geq 3.0$. This gets confirmed by an inspection of Fig. 4. This can be understood from the fact that any enhanced fermionic contribution due to a higher $\tan\beta$ can always be cancelled by an appropriately increased bosonic contribution through a slight tweak in the masses. Of course, one also has to keep the couplings perturbative in doing

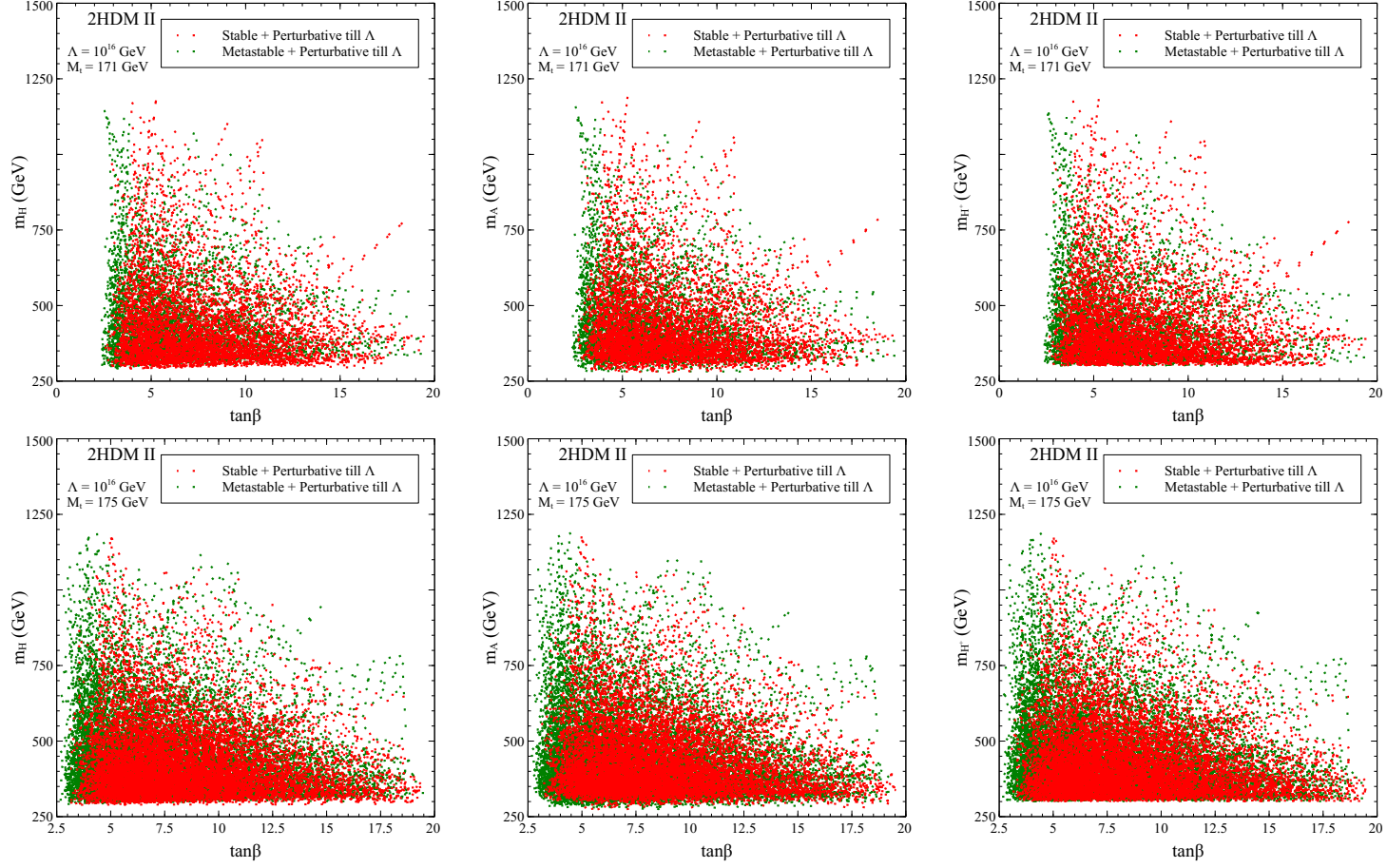


Figure 3: Distribution of models perturbative till 10^{16} GeV that lead to stable, or metastable EW vacuum. The upper (lower) plots correspond to $M_t = 171$ (175) GeV. The color coding is explained in the legends. 2HDM II refers to a Type II 2HDM.

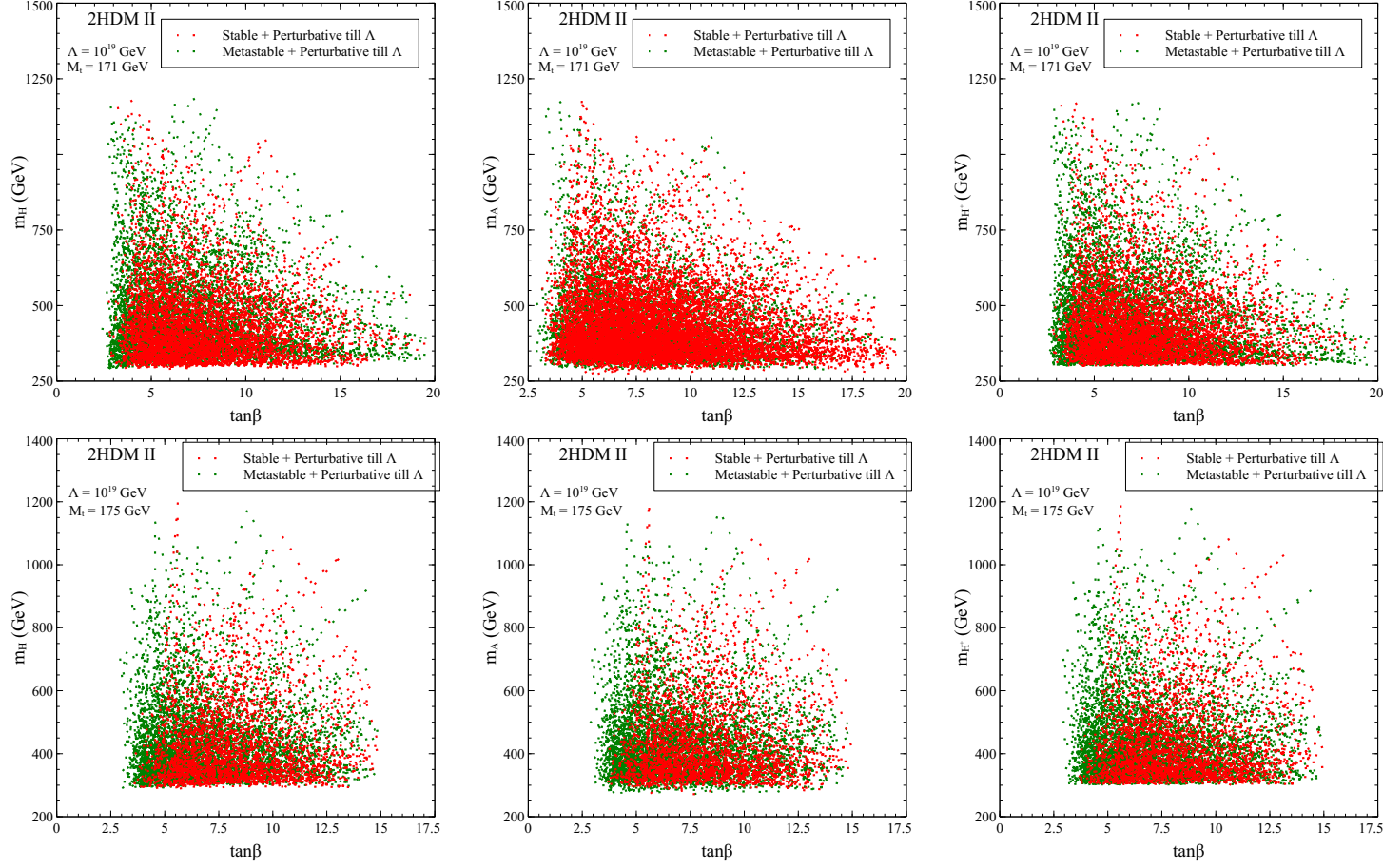


Figure 4: Distribution of models perturbative till 10^{19} GeV that lead to stable, or metastable EW vacuum. The upper (lower) plots correspond to $M_t = 171$ (175) GeV. The color coding is explained in the legends. 2HDM II refers to a Type II 2HDM.

so. Such a "fine-tuned" existence of a metastable EW vacuum is not a surprise and is always expected in the case of an extended Higgs sector, such as the 2HDM.

We take another approach where different scalar masses are fixed within specific narrow ranges, and allow $\tan\beta$ to vary. This approach turns useful in demarcating the "stable" region from the "metastable". We thus propose two central values of 500 GeV and 1000 GeV and allow only a 2 GeV split about that. Fig. 5 presents the results for this choice.

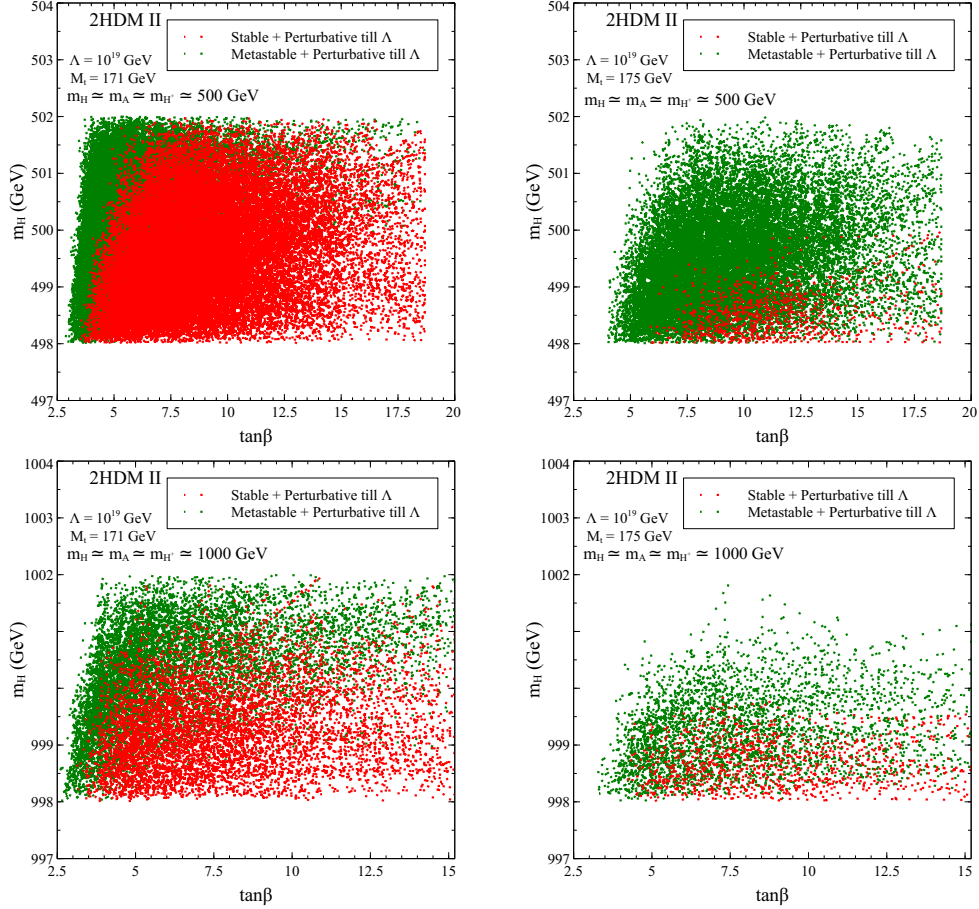


Figure 5: Distribution of models perturbative till 10^{19} GeV that lead to stable, or metastable EW vacuum. The upper (lower) plots correspond to $M_t = 171$ (175) GeV. The color coding is explained in the legends. 2HDM II refers to a Type II 2HDM.

For masses around 500 GeV and $M_t = 171$ GeV, the metastable points mostly cluster in the low $\tan\beta$ region. They get largely disfavored at larger $\tan\beta$. Since the bosonic contribution to RG evolution is now restrained, absolute stability demands $\tan\beta \geq 3.0$. For $M_t = 175$ GeV however, lower bound on $\tan\beta$ for both stability as well as metastability goes up, stability completely ruled out for $\tan\beta \leq 5.0$ for instance. Thus, for $M_t = 175$ GeV, the proportion of metastable model points is higher compared to what is seen for $M_t = 171$ GeV. The robustness of this claim is verified by the plots for masses $\simeq 1000$ GeV, which depict the same qualitative behaviour. Having pointed out the crucial role played by the parameter $\tan\beta$, we close

this section here.

6 Summary and outlook

This work highlights the possibility of a metastable EW vacuum in a popular 2HDM framework. While the occurrence of a panic vacuum in a 2HDM is by and large disfavored by the latest data from the LHC, we observe that a global minimum at high scales is indeed possible, iff RG effects are incorporated into the scalar potential. This is found to happen in the direction of the scalar field h_2 , because λ_2 can be driven to negative values at high scales. We have reported our findings in context of a type II 2HDM.

We remark that it is the relative strengths of the fermionic and bosonic contributions in the RG improved potential that seals the fate of the EW vacuum where we currently reside. The introduction of additional bosonic degrees of freedom further introduces a tension between vacuum stability on the one hand, and, high-scale perturbativity on the other. This tension can be responsible for substantial constraints on the parameter space.

In a 2HDM, the strength of the fermionic contribution is controlled by not only the top quark pole mass, but also $\tan\beta$. Based on the results of this work, one would always expect a metastable model point in the vicinity of a point allowing for absolute stability. However, $\tan\beta$ picks up a lower bound from the requirement of metastability, which is tightened when one demands absolute stability of the EW vacuum. The sensitivity of the results to the top pole mass has also been emphasized.

A pertinent extension would be to include finite temperature corrections to the 2HDM scalar potential and, study its impact on vacuum stability.

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